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Supersymmetric extra U(1) models with a singlino dominated LSP

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Abstract

We investigate phenomenology related to the neutral fields in supersymmetric models with an extra U(1) derived from E_6. Our study is concentrated into the models which have a singlino dominated neutralino as the lightest superparticle (LSP). If such models satisfy a constraint for dark matter derived from the WMAP data, the lightest neutral Higgs scalar, a new neutral gauge field Z' and the LSP may be interesting targets for the study at the LHC. We also discuss features of the Z' in the models and its detectability at the LHC.

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1 Introduction

Recent various astrophysical observations quantitatively show the existence of a substantial amount of non-relativistic and non-baryonic dark matter [1, 2]. Although supersymmetric models have been considered to be the best candidate beyond the standard model (SM) from a viewpoint of both the gauge hierarchy problem and the gauge coupling unification, this fact seems to make them much more promising on the basis of an experimental signature [3]. If R-parity is conserved in supersymmetric models, the lightest superparticle (LSP) is stable. Thus the LSP can be a good candidate for cold dark matter (CDM) as long as it is electrically neutral. Since the strength of interactions of the LSP with the SM fields is $O(G_F)$ and the mass can be of the order of the weak scale, its relic energy density is expected to be eventually of the order of critical energy density of the universe. The most promising one among such LSP candidates is considered to be the lightest neutralino. Relic abundance of the lightest neutralino has been extensively studied in the minimal supersymmetric SM (MSSM) [4, 5, 6]. After publication of the analysis of the WMAP data, however, the allowed parameter space in the minimal SUGRA (or CMSSM) is found to be restricted into some narrow regions [7]. If we take account of this situation, it seems worth studying the relic abundance of CDM candidates quantitatively also in various extensions of the MSSM. It may also be useful to discuss indications of such models which are expected to be found at the forthcoming Large Hadron Collider (LHC) by applying the CDM condition.

The MSSM has been considered as the most promising supersymmetric model and has extensively studied from various points of view. Although the MSSM can explain experimental results obtained by now as long as parameters are suitably chosen, it suffers from the well known μ problem [8]. If we try to solve it near the weak scale, we need to extend the MSSM in the way to give some influence to physics at TeV regions [9, 10]. If we extend it by introducing an extra U(1) gauge symmetry, for example, the problem can be solved in a very elegant way [10, 11]. The existence of additional U(1) symmetries is also predicted in some effective theories of superstring [12, 13]. If there is an extra U(1) symmetry at TeV regions, the model is expected to reveal distinguishable features from those of the MSSM. Some signals of the models may be detected at the LHC [13, 14, 15] and a CDM candidate may be different from that in other models. In this paper we focus our attention into such aspects in the models with an extra U(1).

In the models with an extra U(1), an operator $\lambda\hat{S}\hat{H}_1\hat{H}_2$ is introduced into superpotential as a gauge invariant operator instead of the so-called μ term $\mu\hat{H}_1\hat{H}_2$ [10].[1] In the simplest models to accommodate such a feature, the extra U(1) symmetry is supposed to be broken by a vacuum expectation value (VEV) of the scalar component S of an SM singlet chiral superfield $\hat{S}$. The μ term is generated as $\mu = \lambda\langle S\rangle$ by the same singlet scalar field through the introduced operator. These models show a difference from the MSSM in the neural Higgs sector in addition to the existence of a new neutral gauge field Z' [10, 16]. Since the neutralino sector is extended from the MSSM by a fermionic component $\tilde{S}$ and an extra U(1) gaugino $\tilde{\lambda}_x$, the feature of neutralinos can also be different from that in the MSSM. Various phenomena are influenced by this change [17].

In particular, if the singlino $\tilde{S}$ can dominate the lightest neutralino, distinguishable features from the MSSM are expected to appear in the phenomena relevant to the neutralinos. When $\langle S\rangle$ takes a value of the weak scale, the singlino domination of the lightest neutralino is naturally expected to occur. This is the case in the well known next MSSM (NMSSM) [18] and its modified model (nMSSM) [19]. In the models with an extra U(1), however, there are severe constraints on $\langle S\rangle$ from both mass bounds of Z' which result from the direct search of Z' [20] and bounds on the mixing between Z' and the ordinary Z which result from the electroweak precision measurements [13, 21, 22]. These constraints tend to require that $\langle S\rangle$ should be more than $O(1)$ TeV as long as we do not consider a special situation.[2] Thus, in the simple models with an extra U(1) which is called the UMSSM in [24], it seems unable to expect substantial differences in the lightest neutralino from the MSSM since both $\tilde{S}$ and $\tilde{\lambda}_x$ tend to decouple from the lightest neutralino.

This situation changes if the extra U(1) gaugino $\tilde{\lambda}_x$ is sufficiently heavier than Z' as suggested in [11, 23, 24, 25]. In these papers the nature of the lightest neutralino has been studied to show a lot of interesting aspects: (a) it can be dominated by the singlino $\tilde{S}$ and can be very light; (b) it can be a good candidate of the CDM whose nature is

[1] In this paper we put a hat on the character for a superfield. For its component fields, we put a tilde on the same character to represent the superpartners of the SM fields and use just the same character without the hat for the SM fields. Otherwise, the field with no tilde should be understood as a scalar component.

[2] Even in the models with an extra U(1), if one considers a model with a secluded singlet sector which is called the S-model in [24], $\langle S\rangle$ can take a value of the weak scale. In this case phenomenological feature at the weak scale is very similar to the nMSSM.

very different from that in the MSSM. It can have a small mass which has already been forbidden in the MSSM since it is dominated by the singlino; (c) it also has a different interaction from that in the MSSM. However, the studied parameter regions are different in each study. In [23] the CDM abundance is mainly studied under the assumption that the lightest neutral Higgs mass is $m_h = 170$ GeV. In the S-model case the neutralino mass matrix is assumed to be reduced to the one of the nMSSM [24] or a certain gaugino mass relation such as $M_{\tilde{x}} = M_{\tilde{Y}}$ is assumed [26]. In this paper we do not use these assumptions. We are interested in other region of the parameter space, where $\tilde{\lambda}_x$ is much heavier than other gauginos but it has non-negligible mixing with $\tilde{S}$ [25, 27]. Therefore $\tilde{\lambda}_x$ can give a crucial contribution to the mass of $\tilde{S}$. We focus our study on such cases and investigate the neutral field sector in the models with an extra U(1) derived from E_6. We discuss phenomenological features of the models and study indications of the models expected to be found at the LHC.

The paper is organized as follows. In section 2 we define the models and discuss their features different from other models. In particular, we focus our attention on the neutral field sectors of the models, that is, the lightest neutralino, the new neutral gauge field Z' and the neutral Higgs scalars. In section 3 we study the parameter space of the models which is allowed by the current results of various experiments. We predict signals of the models expected to be seen at the LHC. Section 4 is devoted to the summary.

2 μ problem solvable extra U(1) models

2.1 Features of the neutral fields

In the simple models with an extra U(1) which can solve the μ problem, the mass of Z' is directly related to the μ term [10, 11]. This feature induces various interesting phenomena which make the models distinguishable from the MSSM. The μ term is considered to be generated by an operator in the last term of the superpotential

$$W_{\rm ob} = h_U \hat{\bar{U}}\hat{Q}\hat{H}_2 + h_D \hat{\bar{D}}\hat{Q}\hat{H}_1 + h_E \hat{\bar{E}}\hat{L}\hat{H}_1 + \lambda \hat{S}\hat{H}_1\hat{H}_2, \tag{1}$$

where $\hat{H}_{1,2}$ are the ordinary doublet Higgs chiral superfields and $\hat{S}$ is an additional singlet chiral superfield. If a VEV of the scalar component of $\hat{S}$ is assumed to generate both the μ term and the Z' mass, the superpotential (1) requires that $\hat{H}_{1,2}$ also have the extra

U(1) charges $Q_{1,2}$. The charge conservation imposes a condition $Q_1 + Q_2 + Q_S = 0$ on them. A bare μ term $\mu \hat{H}_1 \hat{H}_2$ is automatically forbidden by this symmetry. Stability of the scalar potential for S is automatically guaranteed by a quartic term induced as the extra U(1) D-term without introducing a new term in the superpotential. The Z' mass is difficult to be much larger than $O(1)$ TeV as long as there is no other contribution to it. These aspects may make the models not only theoretically interesting but also a good target for the studies at the LHC [13, 14, 15]. In particular, we can find interesting features in the neutral fields, which make the models distinguishable from the MSSM and also the singlet extensions of the MSSM such as the NMSSM, the nMSSM and the S-model. In this section we review these features to clarify the differences of our models from previously studied ones.

Before proceeding to the discussion on this issue, we fix assumptions for the supersymmetry breaking. We assume soft supersymmetry breaking terms

$$\begin{aligned} -\mathcal{L}_{\rm SUSY} &= \sum_\varphi m_\varphi^2 |\varphi|^2 + \left(\frac{1}{2} M_{\tilde{g}} \tilde{\lambda}_g \tilde{\lambda}_g + \frac{1}{2} M_{\tilde{W}} \tilde{\lambda}_W \tilde{\lambda}_W + \frac{1}{2} M_{\tilde{Y}} \tilde{\lambda}_Y \tilde{\lambda}_Y + \frac{1}{2} M_{\tilde{x}} \tilde{\lambda}_x \tilde{\lambda}_x + \text{h.c.} \right) \\ &- \left(A_U h_U \hat{\bar{U}} \hat{Q} \hat{H}_2 + A_D h_D \hat{\bar{D}} \hat{Q} \hat{H}_1 + A_E h_E \hat{\bar{E}} \hat{L} \hat{H}_1 + A_\lambda \lambda \hat{S} \hat{H}_1 \hat{H}_2 + \text{h.c.} \right), \end{aligned} \tag{2}$$

where φ in the first term runs all scalar fields contained in the models. Scalar mass m_φ and A parameters for scalar trilinear terms may be assumed to have a universal value $m_{3/2}$. A grand unification relation for the masses of gauginos may also be assumed.

Now we assume that the scalar components of $\hat{H}_{1,2}$ and $\hat{S}$ obtain the VEVs v_1, v_2 and u due to radiative effects on the supersymmetry breaking parameters and the electroweak symmetry $\text{SU}(2)_L \times \text{U}(1)_Y \times \text{U}(1)_x$ breaks down into $\text{U}(1)_{\rm em}$ [10, 11]. We define the field fluctuations around this vacuum as

$$H_1 = \begin{pmatrix} v_1 + h_1^0 + iP_1 \\ h_1^- \end{pmatrix}, \quad H_2 = \begin{pmatrix} h_2^+ \\ v_2 + h_2^0 + iP_2 \end{pmatrix}, \quad S = u + h_S^0 + iP_S. \tag{3}$$

The last term in eq. (1) generates a μ parameter as $\mu = \lambda u$ and plays a required role for the μ term in the MSSM. Although the models have similar structures to the MSSM after this symmetry breaking, there appear various differences induced by this term, in particular, in the sector of the neutral fields. In a charged field sector, we can find a difference from the MSSM only in the mass of the charged Higgs scalar. It will be discussed in Appendix.

2.1.1 Neutral gauge field sector

The clearest difference from the MSSM is the existence of the Z' at TeV regions. Through the symmetry breaking denoted by eq. (3), a mass matrix $M^2_{ZZ'}$ is generated for the neutral gauge bosons Z_μ and Z'_μ. It can be written in the basis (Z_μ, Z'_μ) as

$$\begin{pmatrix} \frac{g_2^2+g_1^2}{2}v^2 & \frac{g_x\sqrt{g_2^2+g_1^2}}{2}v^2(Q_1\cos^2\beta - Q_2\sin^2\beta) \\ \frac{g_x\sqrt{g_2^2+g_1^2}}{2}v^2(Q_1\cos^2\beta - Q_2\sin^2\beta) & \frac{g_x^2}{2}v^2(Q_1^2\cos^2\beta + Q_2^2\sin^2\beta + Q_S^2\frac{u^2}{v^2}) \end{pmatrix} \tag{4}$$

where $v^2 = v_1^2 + v_2^2$ and $\tan\beta = v_2/v_1$. The extra U(1) charge $Q_x(f)$ of a field f and its coupling g_x are defined in such a way that the covariant derivative for the relevant gauge group $\mathrm{SU}(2)_L\times\mathrm{U}(1)_Y\times\mathrm{U}(1)_x$ takes the form

$$D_\mu = \partial_\mu - ig_2\frac{\tau^a}{2}W^a_\mu - ig_1\frac{Y(f)}{2}B_\mu - ig_x\frac{Q_x(f)}{2}Z'_\mu. \tag{5}$$

Mass eigenvalues of the neutral gauge fields $Z_{1\mu}$ and $Z_{2\mu}$ can be approximately expressed as

$$\begin{aligned} m^2_{Z_1} &\simeq m_Z^2 - m_Z^2\frac{g_x\tan 2\xi}{\sqrt{g_2^2+g_1^2}}(Q_1\cos^2\beta - Q_2\sin^2\beta), \\ m^2_{Z_2} &\simeq \frac{g_x^2}{2}(Q_1^2v_1^2 + Q_2^2v_2^2 + Q_S^2u^2) + m_Z^2\frac{g_x\tan 2\xi}{\sqrt{g_2^2+g_1^2}}(Q_1\cos^2\beta - Q_2\sin^2\beta), \end{aligned} \tag{6}$$

where m_Z is the mass of the Z boson in the SM and ξ is a ZZ' mixing angle defined by

$$\tan 2\xi = \frac{2g_x\sqrt{g_2^2+g_1^2}(Q_1\cos^2\beta - Q_2\sin^2\beta)}{g_x^2(Q_1^2\cos^2\beta + Q_2^2\sin^2\beta + Q_S^2u^2/v^2) - (g_2^2+g_1^2)}. \tag{7}$$

If $u \gg v_{1,2}$ is satisfied, m_{Z_1} approaches to m_Z and m_{Z_2} is proportional to u. The relation between m_{Z_2} and u is given by $m_{Z_2}/u \simeq g_xQ_S/\sqrt{2}$.

Both direct search of a new neutral gauge field and precise measurements of the electroweak interaction severely constrain the mass eigenvalue m_{Z_2} of the new neutral gauge boson and the ZZ' mixing angle ξ [21]. Lower bounds for m_{Z_2} have been studied by using the searches of the Z_2 decay into dilepton pairs [20]. Although it depends on the models, it may be roughly estimated as $m_{Z_2} \gtrsim 600$ GeV. If Z_2 has a substantial decay width into non-SM fermion pairs such as neutralino pairs, this bound may be largely relaxed [15]. On the other hand, the precise measurements of the electroweak interaction give a constraint $\xi \lesssim 10^{-3}$ [21]. As found from eq. (7), this bound can be fulfilled if either of two conditions is satisfied, that is, a sufficiently large u or $\tan\beta \simeq \sqrt{Q_1/Q_2}$ [11]. For the latter case, since the constraint from the ZZ' mixing can automatically be guaranteed, u needs not so large as long as the direct search constraint on m_{Z_2} is satisfied. This seems to be an important point to be noted when we consider the existence of an extra U(1) at TeV regions. We will focus our study on this case, which may make it possible to find the solutions in the parameter regions excluded in the study of the UMSSM [23, 24].

2.1.2 Neutralino sector

The neutralino sector is extended into six components, since there are two additional neutral fermions $\tilde{\lambda}_x$ and $\tilde{S}$. $\tilde{\lambda}_x$ is the extra U(1) gaugino and $\tilde{S}$ is the fermionic component of $\hat{S}$. If we take a basis $\mathcal{N}^T = (-i\tilde{\lambda}_x, -i\tilde{\lambda}^3_W, -i\tilde{\lambda}_Y, \tilde{H}_1, \tilde{H}_2, \tilde{S})$ and define neutralino mass terms such as $\mathcal{L}^m_{\rm neutralino} = -\frac{1}{2}\mathcal{N}^T\mathcal{M}\mathcal{N} + {\rm h.c.}$, a 6 $\times$ 6 neutralino mass matrix $\mathcal{M}$ can be represented as[3]

$$\begin{pmatrix} M_{\tilde{x}} & 0 & 0 & \frac{g_xQ_1}{\sqrt{2}}v\cos\beta & \frac{g_xQ_2}{\sqrt{2}}v\sin\beta & \frac{g_xQ_S}{\sqrt{2}}u \\ 0 & M_{\tilde{W}} & 0 & m_Zc_W\cos\beta & -m_Zc_W\sin\beta & 0 \\ 0 & 0 & M_{\tilde{Y}} & -m_Zs_W\cos\beta & m_Zs_W\sin\beta & 0 \\ \frac{g_xQ_1}{\sqrt{2}}v\cos\beta & m_Zc_W\cos\beta & -m_Zs_W\cos\beta & 0 & \lambda u & \lambda v\sin\beta \\ \frac{g_xQ_2}{\sqrt{2}}v\sin\beta & -m_Zc_W\sin\beta & m_Zs_W\sin\beta & \lambda u & 0 & \lambda v\cos\beta \\ \frac{g_xQ_S}{\sqrt{2}}u & 0 & 0 & \lambda v\sin\beta & \lambda v\cos\beta & 0 \end{pmatrix}. \tag{8}$$

Neutralino mass eigenstates $\tilde{\chi}^0_a (a = 1 \sim 6)$ are related to $\mathcal{N}_j$ by using the mixing matrix U as

$$\tilde{\chi}^0_a = \sum_{j=1}^{6} U_{aj}\mathcal{N}_j, \tag{9}$$

where U is defined in such a way that $U\mathcal{M}U^T$ becomes diagonal.

The composition of the lightest neutralino is important for the study of various phenomena, in particular, the relic density of the lightest neutralino as a CDM candidate. If u is a similar order value to $v_{1,2}$ or less than these, the lightest neutralino is expected to be dominated by the singlino $\tilde{S}$ just like in the case of the nMSSM and the S-model. In this case, if it can annihilate sufficiently, the lightest neutralino with a sizable singlino

[3]We do not consider gauge kinetic term mixing between $\mathrm{U}(1)_Y$ and the extra U(1), for simplicity. The study of their phenomenological effects can be found in [17].

component may be a good CDM candidate in the parameter regions different from the ones in the MSSM [18, 19]. In the present models, however, the Z' constraints seem to require that u should be much larger than $v_{1,2}$ as mentioned before. As the result, $\tilde{\lambda}_x$ and $\tilde{S}$ tend to decouple from the lightest neutralino as long as the mass of $\tilde{\lambda}_x$ is assumed to be similar to the masses of other gauginos. In such a situation, the composition of the lightest neutralino is expected to be similar to that of the MSSM. Then we cannot find distinctive features in the lightest neutralino. However, as suggested in [11, 23, 24, 25], this situation can be drastically changed if the mass of the extra U(1) gaugino $M_{\tilde{x}}$ becomes much larger than the masses of other gauginos due to some reasons. In this case the lightest neutralino can be dominated by the singlino $\tilde{S}$.

If the gaugino $\tilde{\lambda}_x$ is heavy enough to satisfy $M_{\tilde{x}} \gg \frac{g_x Q_S}{\sqrt{2}} u$, we can integrate out $\tilde{\lambda}_x$ just as the seesaw mechanism for the neutrinos. A resulting 5×5 mass matrix can be expressed as

$$\begin{pmatrix} M_{\tilde{W}} & 0 & m_Z c_W \cos\beta & -m_Z c_W \sin\beta & 0 \\ 0 & M_{\tilde{Y}} & -m_Z s_W \cos\beta & m_Z s_W \sin\beta & 0 \\ m_Z c_W \cos\beta & -m_Z s_W \cos\beta & -\frac{g_x^2 Q_1^2}{2M_{\tilde{x}}} v^2 \cos^2\beta & \lambda u & \lambda v \sin\beta \\ -m_Z c_W \sin\beta & m_Z s_W \sin\beta & \lambda u & -\frac{g_x^2 Q_2^2}{2M_{\tilde{x}}} v^2 \sin^2\beta & \lambda v \cos\beta \\ 0 & 0 & \lambda v \sin\beta & \lambda v \cos\beta & -\frac{g_x^2 Q_S^2}{2M_{\tilde{x}}} u^2 \end{pmatrix}. \tag{10}$$

This effective mass matrix suggests that the lightest neutralino tends to be dominated by the singlino $\tilde{S}$ as long as $M_{\tilde{W},\tilde{Y}}$ and $\mu(\equiv \lambda u)$ is not smaller than $\frac{g_x^2 Q_S^2 u^2}{2M_{\tilde{x}}}$. Since $M_{\tilde{W}}$ and μ cannot to be less than 100 GeV because of mass bounds of the lightest chargino and the gluino [28], the singlino domination of the lightest neutralino is expected to be naturally realized in the case that $M_{\tilde{x}} \gg u$ is satisfied. In such a case, phenomenology of the lightest neutralino can be largely changed from that in the MSSM and also its singlet extensions. This effective mass matrix reduces to the one of the nMSSM in the large $M_{\tilde{x}}$ limit [24]. However, we are interested in the intermediate situation where the effectively generated diagonal elements in eq. (10) can not be neglected. Since this possibility has not been studied in detail in realistic models yet despite it is potentially interesting,[4] we will concentrate our study into such a situation in this paper. The lightest neutralino is assumed to be dominated by the singlino because of a large $M_{\tilde{x}}$ compared with the mass

[4]In particular, the detailed study seems to have not been done under the assumption $\tan\beta = \sqrt{Q_1/Q_2}$. In [25] we do not consider the anomaly problem of $\mathrm{U}(1)_x$ seriously and only study a toy model.

of other gauginos.

2.1.3 Neutral Higgs scalar sector

The neutral Higgs mass is also modified from that in the MSSM as in the case of the NMSSM and the nMSSM. A difference from the latter ones is the existence of an additional D-term contribution of the extra U(1). The CP even neutral Higgs sector is composed of the three scalars (h_1^0, h_2^0, h_S^0) which are introduced in eq. (3). Their mass matrix $\mathcal{M}_h^2$ at tree level is written as

$$\begin{pmatrix} \frac{1}{2}(g_2^2+g_1^2+g_x^2Q_1^2)v_1^2 + A_\lambda\lambda u\tan\beta & -\frac{1}{2}(g_2^2+g_1^2-g^2(1,2))v_1v_2 - A_\lambda\lambda u & \frac{1}{2}g^2(1,S)v_1u - A_\lambda\lambda v_2 \\ -\frac{1}{2}(g_2^2+g_1^2-g^2(1,2))v_1v_2 - A_\lambda\lambda u & \frac{1}{2}(g_2^2+g_1^2+g_x^2Q_2^2)v_2^2 + A_\lambda\lambda u\cot\beta & \frac{1}{2}g^2(2,S)v_2u - A_\lambda\lambda v_1 \\ \frac{1}{2}g^2(1,S)v_1u - A_\lambda\lambda v_2 & \frac{1}{2}g^2(2,S)v_2u - A_\lambda\lambda v_1 & \frac{1}{2}g_x^2Q_S^2u^2 + A_\lambda\lambda\frac{v_1v_2}{u} \end{pmatrix}, \tag{11}$$

where A_λ is the soft supersymmetry breaking parameter defined in eq. (2). We use a definition $g^2(i,j) = g_x^2Q_iQ_j + 4\lambda^2$ in this formula. Mass eigenstates ϕ_α are related to the original neutral CP even Higgs scalars h_a^0 by

$$\phi_\alpha = \sum_{a=1,2,S} \mathcal{O}_{\alpha a}h_a^0, \tag{12}$$

where the orthogonal matrix $\mathcal{O}$ is defined so as to diagonalize the neutral Higgs mass matrix $\mathcal{M}_h^2$ in such a way as $\mathcal{O}\mathcal{M}_h^2\mathcal{O}^T = \mathrm{diag}(m_{\phi_1}^2, m_{\phi_2}^2, m_{\phi_3}^2)$.

Since upper bounds of the mass eigenvalue for the lightest neutral Higgs scalar h^0 can be estimated by using eq.(11) as

$$m_{h^0}^2 \leq m_Z^2\left[\cos^2 2\beta + \frac{2\lambda^2}{g_2^2+g_1^2}\sin^2 2\beta + \frac{g_x^2}{g_2^2+g_1^2}(Q_1\cos^2\beta + Q_2\sin^2\beta)^2\right] + \Delta m_1^2, \tag{13}$$

it can be larger than that in the MSSM. The second term in the brackets of eq. (13) comes from the interaction given by the last term in eq. (1). It can give a large contribution for smaller values of $\tan\beta$ and u for a fixed μ. The third term is the D-term contribution of the extra U(1). Due to these effects, even in the regions of the small $\tan\beta$ such as $\tan\beta = 1 - 2$, the mass of the lightest neutral Higgs m_{ϕ_α} can take larger values than that in the MSSM, such as 120 GeV or more, if one-loop corrections Δm_1^2 are taken into account [16]. Since dominant components of this lightest Higgs scalar are expected to be $h_{1,2}^0$ as long as u is not smaller than $v_{1,2}$, its nature is similar to that in the MSSM, except that it is heavier than that in the MSSM. This situation seems very different from the nMSSM and the S-model where the lightest neutral Higgs is a mixture state of $h^0_{1,2}$ and h^0_S. In numerical studies given in section 3, we will estimate both the Higgs mass eigenvalues and their eigenstates by diagonalizing the mass matrix (11) including the one-loop corrections due to the stops.

CP odd Higgs scalars are also somewhat changed from the ones in the MSSM. A CP odd Higgs mass matrix $\mathcal{M}_P^2$ can be written as

$$\mathcal{M}_P^2 = \begin{pmatrix} A_\lambda\lambda u\tan\beta & A_\lambda\lambda u & A_\lambda\lambda v_2 \\ A_\lambda\lambda u & A_\lambda\lambda u\cot\beta & A_\lambda\lambda v_1 \\ A_\lambda\lambda v_2 & A_\lambda\lambda v_1 & A_\lambda\lambda\frac{v_1v_2}{u} \end{pmatrix}. \tag{14}$$

Only one component P_A has a non-zero mass eigenvalue

$$m_{P_A}^2 = \frac{2A_\lambda\lambda u}{\sin 2\beta}\left(1+\frac{v^2}{4u^2}\sin^2 2\beta\right), \tag{15}$$

and others are would-be Goldstone bosons $G^0_{1,2}$ as in the MSSM. This requires $\lambda u A_\lambda > 0$ for the stability of the vacuum. Imaginary parts of the original Higgs fields in eq. (3) have P_A as a component. They can be written as

$$P_1 = \frac{u\sin\beta}{N}P_A + \ldots, \quad P_2 = \frac{u\cos\beta}{N}P_A + \ldots, \quad P_S = \frac{v\sin\beta\cos\beta}{N}P_A + \ldots, \tag{16}$$

where a normalization factor N is defined as $N=\sqrt{u^2+v^2\sin^2\beta\cos^2\beta}$. Although $m^2_{P_A}$ takes larger values than those in the MSSM, P_A is found to be similar to that in the MSSM if u becomes larger than v.

2.2 A CDM constraint

In the models with a large $M_{\tilde{x}}$ the composition and the interaction of the lightest neutralino can be very different from that in the MSSM and its singlet extensions. Since the lightest neutralino can also be a CDM candidate in such models, the relic abundance is expected to give a different constraint on the parameter space from that in the MSSM and its singlet extensions. It is useful to discuss this point briefly here.

The relic abundance of the stable lightest neutralino $\tilde{\chi}^0_\ell$ which is thermally produced can be evaluated as thermal abundance at its freeze-out temperature T_F. This temperature can be determined by $H(T_F)\sim\Gamma_{\tilde{\chi}^0_\ell}$. $H(T_F)$ is the Hubble parameter at T_F [29]. $\Gamma_{\tilde{\chi}^0_\ell}$ is an annihilation rate of $\tilde{\chi}^0_\ell$ and it can be expressed as $\Gamma_{\tilde{\chi}^0_\ell}=\langle\sigma_{\rm ann}v\rangle n_{\tilde{\chi}^0_\ell}$, where $\langle\sigma_{\rm ann}v\rangle$ is thermal average of the product of an annihilation cross section $\sigma_{\rm ann}$ and relative velocity v of annihilating $\tilde\chi_\ell^0$s. Thermal number density of non-relativistic $\tilde\chi_\ell^0$ at this temperature is expressed by $n_{\tilde\chi_\ell^0}$. If we use the non-relativistic expansion for the annihilation cross section such as $\sigma_{\rm ann}v \simeq a + bv^2$ and introduce a dimensionless parameter $x_F = m_{\tilde\chi_\ell^0}/T_F$, we find that x_F can be represented as

$$x_F = \ln \frac{0.0955 m_{\rm pl} m_{\tilde\chi_\ell^0}(a + 6b/x_F)}{(g_* x_F)^{1/2}}, \tag{17}$$

where $m_{\rm pl}$ is the Planck mass and g_* enumerates the degrees of freedom of relativistic particles at T_F. Using this x_F, the present abundance of $\tilde\chi_\ell^0$ can be estimated as

$$\Omega_\chi h^2|_0 = \left.\frac{m_{\tilde\chi_\ell^0} n_{\tilde\chi_\ell^0}}{\rho_{\rm cr}/h^2}\right|_0 \simeq \frac{8.76\times 10^{-11} g_*^{-1/2} x_F}{(a + 3b/x_F)\ {\rm GeV}^2}. \tag{18}$$

We can find formulas of the coefficients a and b for the processes mediated by the exchange of various fields contained in the MSSM in the articles [4, 5].

If the lightest neutralino $\tilde\chi_\ell^0$ is dominated by the singlino and also relatively light, the decay modes into other final states than the SM fermion-antifermion pairs are expected to be suppressed. The singlino dominated neutralino contains the MSSM higgsinos ($\mathcal{N}_{4,5}$) and gauginos ($\mathcal{N}_{2,3}$) as its components with an extremely small ratio. Thus the annihilation process caused by these components are heavily suppressed and then cannot be dominant modes unless the enhancement due to the pole effects of the intermidiate fields. As long as we consider this lightest neutralino is relatively light, such pole enhancements are kinematically forbidden in the annihilation modes which have gauge bosons (W^+W^-, Z_1Z_1) and Higgs scalars as the final states. Then we can expect that the annihilation through the modes $\tilde\chi_\ell^0\tilde\chi_\ell^0 \to f\bar f$ is dominant, since the above mentioned situation is escapable for this process. Moreover, exotic fields are considered to be heavy enough and then cannot be the final states kinematically. Thus f is expected to be restricted to quarks and leptons. In the present analyses we will mainly focus our attention to this case. These annihilation processes of the lightest neutralino in the models with an extra U(1) are expected to be mediated by the exchange of Z_1, Z_2 and the neutral Higgs scalars in the s-channel and by the sfermion exchange in the t-channel. New interactions related to these annihilation processes of the lightest neutralino $\tilde\chi_\ell^0$ can be written as

$$\mathcal{L} = \sum_{j=4}^{6} \frac{g_x Q_j}{2}\bar{\mathcal{N}}_j\gamma_5\gamma^\mu\mathcal{N}_j Z'_\mu + \lambda\left(h_1^0\mathcal{N}_5\mathcal{N}_6 + h_2^0\mathcal{N}_4\mathcal{N}_6 + h_3^0\mathcal{N}_4\mathcal{N}_5\right),$$

$$
\begin{aligned}
&+ \quad \frac{g_x Q(f)}{\sqrt{2}}\left(\bar{\mathcal{N}}_1 \bar{f}\tilde{f} - \mathcal{N}_1 f \tilde{f}^*\right) + \dots, \\
&\simeq \quad \sum_{j=4}^{6} \frac{g_x Q_j}{2} U_{\ell j}^2 \bar{\tilde{\chi}}_j^0 \gamma_5 \gamma_\mu \tilde{\chi}_j^0 Z_2^\mu + \lambda U_{\ell 6}(U_{\ell 4}\mathcal{O}_{\alpha 2} + U_{\ell 5}\mathcal{O}_{\alpha 1})\tilde{\chi}_\ell^0 \tilde{\chi}_\ell^0 \phi_\alpha \\
&+ \quad \frac{g_x Q(f)}{\sqrt{2}} U_{\ell 1}\tilde{\chi}_\ell^0 \left(\bar{f}\tilde{f} - f\tilde{f}^*\right) + \dots,
\end{aligned}
\tag{19}
$$

where we use eqs. (9) and (12).

If we consider the case that the singlino dominates the lightest neutralino, the annihilation cross section into the final states $f\bar{f}$ is expected to obtain the dominant contributions from the exchange of the new neutral gauge field Z_2 and the exchange of the lightest neutral Higgs scalar ϕ_α. They crucially depend on both the composition and the mass of $\tilde{\chi}_\ell^0$ and ϕ_α. These contributions to a and b can be expressed as [23]

$$
\begin{aligned}
a_f &= \frac{2c_f}{\pi}\left[\frac{m_f g_x^2 \sum_{j=4}^{6}\frac{Q_j}{2}U_{\ell j}^2}{4m_{\tilde{\chi}_\ell^0}^2 - m_{Z_2}^2}\left(\frac{Q(f_L)}{2} - \frac{Q(f_R)}{2}\right)\right]^2 \left(1 - \frac{m_f^2}{m_{\tilde{\chi}_\ell^0}^2}\right)^{1/2} + \dots, \\
b_f &= \frac{1}{6}\left(-\frac{9}{2} + \frac{3}{4}\frac{m_f^2}{m_{\tilde{\chi}_\ell^0}^2 - m_f^2}\right) a_f \\
&+ \frac{c_f}{3\pi}\left[\frac{m_{\tilde{\chi}_\ell^0} g_x^2 \sum_{j=4}^{6}\frac{Q_j}{2}U_{\ell j}^2}{4m_{\tilde{\chi}_\ell^0}^2 - m_{Z_2}^2}\right]^2 \left[\left(\frac{Q(f_L)}{2}\right)^2 + \left(\frac{Q(f_R)}{2}\right)^2\right]\left(4 + \frac{2m_f^2}{m_{\tilde{\chi}_\ell^0}^2}\right)\left(1 - \frac{m_f^2}{m_{\tilde{\chi}_\ell^0}^2}\right)^{1/2} \\
&+ \frac{c_f}{8\pi}\left(\frac{m_{\tilde{\chi}_\ell^0}}{4m_{\tilde{\chi}_\ell^0}^2 - m_{\phi_\alpha}^2}\frac{\lambda m_f}{v}\mathcal{P}_f\right)^2 \left(1 - \frac{m_f^2}{m_{\tilde{\chi}_\ell^0}^2}\right)^{3/2} + \dots,
\end{aligned}
\tag{20}
$$

where $c_f = 1$ for leptons and 3 for quarks. The extra U(1) charges of the fermions $f_{L,R}$ are denoted by $Q(f_L)$ and $Q(f_R)$. $\mathcal{P}_f$ is defined by using eq. (16) as

$$
\mathcal{P}_f = \left\{ \begin{array}{ll} \frac{1}{\sin\beta} U_{\ell 6}\mathcal{O}_{\alpha 2}(U_{\ell 4}\mathcal{O}_{\alpha 2} + U_{\ell 5}\mathcal{O}_{\alpha 1}) & (f \text{ with } T_3 = \frac{1}{2}) \\ \frac{1}{\cos\beta} U_{\ell 6}\mathcal{O}_{\alpha 1}(U_{\ell 4}\mathcal{O}_{\alpha 1} + U_{\ell 5}\mathcal{O}_{\alpha 2}) & (f \text{ with } T_3 = -\frac{1}{2}) \end{array}\right. ,
\tag{21}
$$

where T_3 is the weak isospin. Contributions due to the CP odd and heavier CP even Higgs scalars are represented by the ellipses in eq. (20). They are expected to be suppressed because of their large masses.

Since the second term of b_f has no suppression from the masses of the final state fermions, all quarks and leptons can contribute to this term as long as the threshold is opened. This contribution can be effective for a larger $m_{\tilde{\chi}_\ell^0}(< m_{Z_2})$ even in the case that a_f is suppressed by a large value of m_{Z_2}. The third term of b_f is suppressed by the final state fermion mass but it can give a large contribution in the case of $\phi_\alpha \simeq 2m_{\tilde{\chi}^0_\ell}$, as is well known. In the present models the lightest neutralino can be much lighter than that of the MSSM. On the other hand, the lightest CP even neutral Higgs scalar can be heavier than that of the MSSM as discussed in the previous part. These features may make the Higgs-pole enhancement effective in the annihilation of the lightest neutralino. This is not realized in the nMSSM where the lightest neutral Higgs can be very light due to the singlet-doublet mixture. Since the lightest neutralino can have substantial components which have a new interaction with the lightest neutral Higgs scalar as shown in eq. (19), the Higgs exchange is expected to be important if the lightest neutralino contains the ordinary Higgsino or bino component sufficiently. These aspects have been shown partially by using numerical studies in [25]. In addition to these effects we also have to take account of all other processes mediated by the exchange of the MSSM contents in the numerical estimation of the relic abundance of the lightest neutralino because of the following reasons. Firstly, the lightest neutralino with sufficient Higgsino components may be light enough. In that case, if $m_{\tilde{\chi}^0_\ell} \simeq m_{Z_1}/2$ and $|U_{14}|^2 \neq |U_{15}|^2$ are satisfied as it happens in the S-model [24], the annihilation can be enhanced. Secondly, if the D-term contribution of the extra U(1) makes the masses of sfermions small enough, the t-channel exchange of those sfermions can be a crucial process for the annihilation of the lightest neutralinos. The D-term contribution to the sfermion mass is given in eq. (34) of Appendix.

2.3 Z' decay

Search of Z' is one of important subjects planed at the LHC except for the search of Higgs scalars and superpartners [13, 14]. Since the present models are characterized by the existence of both the new neutral gauge boson Z' and the neutral Higgs scalar heavier than that in the MSSM,[5] the combined analyses of these may give a useful clue for the search of this type of models. Here we present some useful formulas for the study of the Z' at the LHC.

[5]It is interesting that in a completely different context there exist other models which predict both a new neutral gauge boson and a neutral Higgs scalar heavier than that expected in the MSSM [30].

A tree level cross section for the process $pp(p\bar{p}) \to Z_2 X \to f\bar{f}X$ is given as [14]

$$\sigma^f = \sum_q \int_0^1 dx_1 \int_0^1 dx_2 \; \sigma(sx_1x_2; q\bar{q} \to f\bar{f}) G_A^q(x_1, x_2, m_{Z_2}^2)\theta(x_1x_2s - M_\Sigma), \tag{22}$$

where $x_{1,2}$ is defined by $x = \sqrt{\frac{Q^2}{s}}e^y$ using the rapidity y and a squared momentum transfer Q^2. The sum of the masses of final state particles is represented by M_Σ. s is a square of the center of mass energy in a collision. A function $G_A^q(x_1, x_2, m_{Z_2}^2)$ depends on the structure functions of quarks. In the present case, eq. (22) can be approximated as [14]

$$\sigma^f = \frac{\kappa}{s}\frac{4\pi^2}{3}\frac{\Gamma_{Z_2}}{m_{Z_2}}B(f\bar{f})\left[B(u\bar{u}) + \frac{1}{C_{ud}}B(d\bar{d})\right] C\exp\left(-\mathcal{A}\frac{M_{Z_2}}{\sqrt{s}}\right), \tag{23}$$

where $C_{ud} = 2(25)$, $C = 600(300)$ and $\mathcal{A} = 32(20)$ for $pp(p\bar{p})$ collisions. The QCD correction is taken into account by κ and it is fixed to be $\kappa \sim 1.3$ in the following numerical calculations. Γ_{Z_2} is a total width of Z_2 and $B(f\bar{f})$ is a branching ratio of the Z_2 decay into $f\bar{f}$. Formulas for possible decay modes of the Z_2 are summarized in Appendix.

We should note that σ^f may be expected to take rather different values from the ordinary ones if the singlino dominated lightest neutralino can explain the observed CDM abundance. Since the decay width into the neutralino sector can be enhanced in comparison with the ordinary Z' models, the detectability of the Z_2 at the LHC may receive a large influence as long as the Z_2 is searched by using dilepton events ($f = e, \mu$) [23, 24]. We will compare it with the results obtained in the ordinary Z' models by practicing numerical analyses in the next section.

3 Numerical analyses

3.1 Set up for the analyses

In this section we study parameter space of the models allowed by various phenomenological constraints including the CDM condition obtained from the analysis of the WMAP data. Then for such parameters we give some predictions of the models for the masses of the new neutral gauge field and the lightest neutral Higgs scalar and the detectability of the Z_2 at the LHC and so on.

Before proceeding to the results of the analyses, we summarize the assumptions which we make in numerical studies. Firstly, we focus our study on the case that the extra U(1)

	Y	Q_ψ	Q_χ
H_1	-1	$-2/\sqrt{6}$	$-2/\sqrt{10}$
H_2	1	$-2/\sqrt{6}$	$2/\sqrt{10}$
Q_L	$1/3$	$1/\sqrt{6}$	$-1/\sqrt{10}$
L_L	-1	$1/\sqrt{6}$	$3/\sqrt{10}$

Table 1 Abelian charges of the relevant fields in a fundamental representation **27** of E_6.

gaugino $\tilde{\lambda}_x$ has a larger mass $M_{\tilde{x}}$ compared with other gauginos. This tends to make the lightest neutralino dominated by the singlino. Secondly, we consider a special case such that $\tan\beta \simeq \sqrt{Q_1/Q_2}$ is satisfied so as to relax the constraint on the value of u and make a phenomenological role of the Z_2 more effective. In this case the ZZ' mixing constraint disappears and only the constraint derived from the direct search of the Z_2 should be taken into account. Thirdly, we restrict our study to the extra U(1)s derived from E_6. The models with an extra U(1) are constrained from anomaly free conditions to be realistic. They are generally required to introduce the exotic matter fields to cancel the anomaly. If we adopt E_6 as a background symmetry, we can control the field contents and their U$(1)_x$ charge systematically to make the model anomaly free. It is also a promising candidate for the extra U(1) models since superstring may realize them as its low energy effective theory [12].

As is well known, E_6 has two Abelian factor groups in addition to the usual SM gauge group. We assume that only one of them remains unbroken at TeV regions and it is broken by the VEV u of the SM singlet scalar given in eq. (3). In this case the general U$(1)_x$ can be expressed as a linear combination of two representative U(1)s such as [14]

$$Q_x = Q_\psi \cos\theta - Q_\chi \sin\theta, \tag{24}$$

where Q_ψ and Q_χ are the charges of U$(1)_\psi$ and U$(1)_\chi$ which are obtained as

$$E_6 \supset SO(10)\times U(1)_\psi, \qquad SO(10)\supset SU(5)\times U(1)_\chi. \tag{25}$$

Each charge of the relevant MSSM fields contained in the fundamental representation **27** of E_6 is given in Table 1. The charge of other chiral superfields in the MSSM can be determined from them by requiring that the superpotential (1) should be invariant under

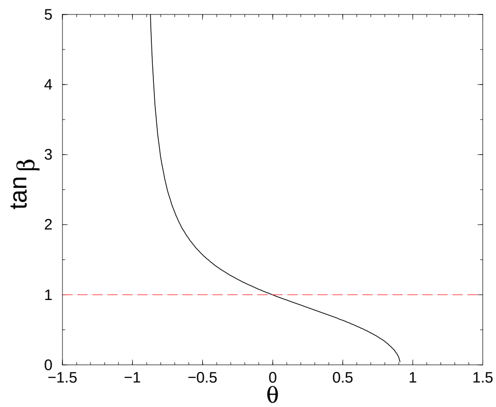

Fig. 1 $\tan\beta$ for various values of θ.

these. Although there are exotic fields such as $\mathbf{3}+\mathbf{3}^*$ of SU(3) and $\mathbf{2}+\mathbf{2}^*$ of SU(2) in $\mathbf{27}$ which are not included in the superpotential (1), they can be assumed to be sufficiently heavy due to some large VEVs or the soft supersymmetry breaking effects. Thus, we neglect these effects on the annihilation of the lightest neutralino and the Z_2 decay in the present analyses. Since we consider that $\mathrm{U}(1)_x$ is derived from E_6, the coupling constant of $\mathrm{U}(1)_x$ may be related to that of the weak hypercharge by $g_x = \sqrt{\frac{5}{3}}g_1$, which is derived from the unification relation

$$\frac{5}{3}g_1^2 \sum_{f\in\mathbf{27}} Y_f^2 = g_x^2 \sum_{f\in\mathbf{27}} Q_x^2.$$

We adopt this relation in the present studies.

In Fig. 1 we plot $\tan\beta$ for the angle θ which is used in eq. (24) to define $\mathrm{U}(1)_x$. Since we consider the case of $\tan\beta = \sqrt{Q_1/Q_2}$, $Q_1/Q_2 \geq 1$ should be satisfied and then the angle θ is found to be confined into the regions such as $-0.9 \lesssim \theta \lesssim 0$. It is interesting that $\theta = -\tan^{-1}(1/\sqrt{15}) \simeq -0.253$ is included in this region. Since the right-handed neutrinos do not have the $\mathrm{U}(1)_x$ charge in this case, they can be very heavy without breaking $\mathrm{U}(1)_x$ and the seesaw mechanism can work to realize the small neutrino mass [10, 22]. We will study this case in detail as an interesting example in the following.

Now we list up free parameters in the models for the numerical analyses. In relation to the extra $\mathrm{U}(1)_x$ we have the VEV u and the angle θ. Since $\tan\beta$ is assumed to be fixed by the $\mathrm{U}(1)_x$ charge as mentioned above, only $\lambda(\equiv \mu/u)$ is a free parameter in the MSSM sector. For the soft supersymmetry breaking parameters, we assume the universality such as $m_\varphi = m_0$ and $A_U = A_D = \cdots = A$ for the parameters in eq. (2). It can be considered as a result of E_6. Thus, if we assume the unification relation for the masses of gauginos $M_{\tilde{g}}$, $M_{\tilde{W}}$, $M_{\tilde{Y}}$ which can be written as $M_{\tilde{g}} = g_3^2 M_{\tilde{W}}/g_2^2$, $M_{\tilde{Y}} = 5g_1^2 M_{\tilde{W}}/3g_2^2$, and also impose $m_0 = A = m_{3/2}$ only to simplify the analyses, the number of remaining parameters is reduced into six:[6]

$$\theta, \quad u, \quad \lambda, \quad M_{\tilde{W}}, \quad M_{\tilde{x}}, \quad m_{3/2}.$$

We practice the following numerical analyses by scanning the parameters $u, M_{\tilde{x}}$ and λ in the following regions:

$$\begin{aligned} &300 \text{ GeV} \leq u \leq 2300 \text{ GeV } (2 \text{ GeV}), \quad 200 \text{ GeV} \leq M_{\tilde{x}} \leq 12 \text{ TeV } (20 \text{ GeV}), \\ &200 \text{ GeV} \leq M_{\tilde{W}}, \ \mu \leq 1300 \text{ GeV } (20 \text{ GeV}), \end{aligned} \tag{26}$$

where it should be noted that μ stands for λ for a fixed u. In the parentheses we show search intervals for these parameters. We fix the supersymmetry breaking parameter $m_{3/2}$ to be 1 TeV as its typical value.[7]

Throughout the analysis we impose the constraints on the masses of the chargino $\tilde{\chi}^{\pm}$, sfermions $\tilde{f}$, the lightest CP even neutral Higgs scalar h^0 and the charged Higgs scalar $h^{\pm}$ as follows:

$$\begin{aligned} m_{\tilde{\chi}^{\pm}} \geq 104 \text{ GeV}, \quad M_{\tilde{g}} \geq 195 \text{ GeV}, \quad m_{\tilde{f}} \geq 250 \text{ GeV}, \\ m_{h^0} \geq 114 \text{ GeV}, \quad m_{h^{\pm}} \geq 79 \text{ GeV}. \end{aligned} \tag{27}$$

Although the constraint on the sfermion masses are model dependent, we use the bound for the constrained MSSM here as an example. The sfermion masses should be checked whether the above bounds are satisfied by including a D-term contribution of the extra U(1). Since the D-term contribution may take a large negative value, this condition could give upper bounds for the value of u in that case. The bounds for m_{Z_2} derived from the direct search of the Z_2 are taken into account by imposing the Tevatron constraint

$$\sigma(p\bar{p} \to Z_2 X) B(Z_2 \to e^+e^-, \ \mu^+\mu^-) < 0.04 \text{ pb} \tag{28}$$

[6]It should be noted that this reduction of the model parameters are caused by the additional assumptions such as $\tan\beta = \sqrt{Q_1/Q_2}$ and $m_0 = A$ which are not related to E_6.

[7]Since $M_{\tilde{x}}$ is assumed to take large values, we need to consider its renormalization group effects on other soft supersymmetry breaking parameters and also the relation to the radiative symmetry breaking. However, since it depends on supersymmetry breaking scenarios, we do not go further this subject here. A relevant study can be found in [27].

at $\sqrt{s} = 1.8$ TeV [20]. We also impose $0 \leq \lambda \equiv \mu/u \leq 0.75$ which is required by perturbative bounds for the coupling constant λ [19, 24, 16]. λ is also restricted indirectly through the relation to μ by the chargino mass bound. The CDM constraint from the 1st year WMAP results is taken into account as [1]

$$0.0945 \leq \Omega_{\rm CDM} h^2 \leq 0.1287 \quad (\text{at } 2\sigma). \tag{29}$$

If there exist other dark matter candidates than the lightest neutralino, the lower bound of (29) cannot be imposed to constrain the parameter region. In this paper, however, we make the analysis under the assumption that the lightest neutralino is the only candidate for the dark matter.

3.2 Predictions of the models

At first, in order to see global features of the extra U(1) models derived from E_6 we study the θ dependence of important quantities. We scan the parameters in the regions shown in (26) for each value of θ which is allowed from Fig. 1 and search the solutions which satisfy all conditions (27)–(29). In Fig. 2 we plot important physical quantities obtained for these solutions. In the left-hand figure we show the predicted regions of the masses of the lightest neutralino $\tilde{\chi}^0_\ell$, the lightest neutral Higgs scalar h^0 and the new neutral gauge boson Z_2. In the right-hand figure we show the predicted cross section for the Z_2 decay into the dilepton final state $(e^+e^-, \mu^+\mu^-)$ at the LHC with $\sqrt{s} = 14$ TeV. The solutions are found only for more restricted values of θ than those shown in Fig. 1. Figure 2 suggests that different annihilation processes of the lightest neutralino play an important role to realize its appropriate relic density for the CDM in different regions of θ. The variety of the predicted neutralino mass shows this aspect.

As interesting results for m_{Z_2}, we find that there exist upper bounds for m_{Z_2}. Since the models considered in the present analysis have small $\tan\beta$ values, the mass bounds of the lightest neutral Higgs scalar can give a constraint on λ as found from eq. (13). This constraint restricts allowed regions of u through the relation $\mu = \lambda u$. On the other hand, we can find the solutions only for restricted regions of μ around 650 – 800 GeV in the present study. The upper bounds of m_{Z_2} appear through these backgrounds. We also find that the lower bound of m_{Z_2} can be less than 600 GeV in the allowed regions of θ. Since the singlino domination of the lightest neutralino makes the decay width of the Z_2

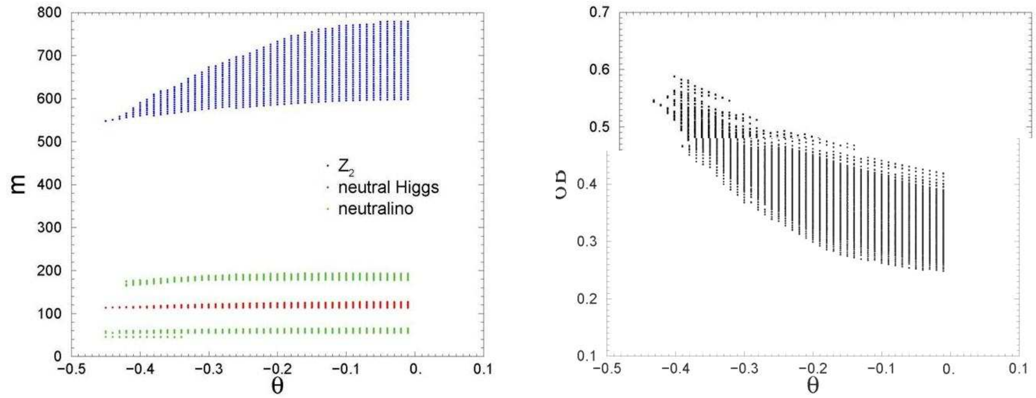

Fig. 2 The left-hand figure shows the masses of the lightest neutral Higgs, the lightest neutralino and the Z_2 boson in a GeV unit for various values of θ. The right-hand figure shows an expected cross section in a fb unit for the Z_2 decay into the dilepton final states at the LHC with $\sqrt{s} = 14$ TeV. These are derived for the parameters which satisfy all conditions discussed in the text.

into the neutralino sector larger, the decay width into the dilepton final state becomes relatively smaller. Thus the constraint imposed by the analysis of the Z_2 search at the CDF seems to be escapable even for a smaller m_{Z_2} than the usually discussed values. The predicted value of σB shows that Z_2 in the present models is easily detectable at the LHC. Combined searches of the Z_2 and the lightest neutral Higgs scalar will be useful to discriminate the models from other candidates beyond the SM.

3.3 Details of annihilation processes of the lightest neutralino

As mentioned in the previous part, different annihilation processes seem to play a crucial role for different regions of θ. To clarify these aspects, we examine the models defined by typical values of θ in detail. We choose $\theta = -0.4$ and -0.253 as such examples. The features found in these examples are also seen in the models defined by other values of θ.

In Fig. 3 we plot points in the $(m_{Z_2}, M_{\tilde{x}})$ plane, which satisfy all of the required conditions in the cases with $\mu = 700$ GeV and representative values of $M_{\tilde{W}}$. The used values of $M_{\tilde{W}}$ are shown in the figures. In each figure we also show the regions surrounded by a dotted line ($\theta = -0.4$) and a dashed line ($\theta = -0.253$) where the lightest neutralino is dominated by the singlino $\tilde{S}$ and also the conditions except for (29) are satisfied. The boundary expressed by a vertical line at a smaller m_{Z_2} is caused by the condition (28)

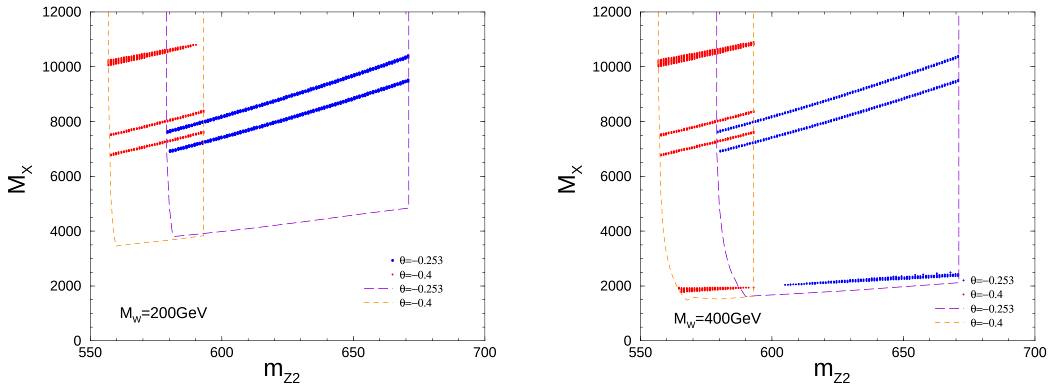

Fig. 3 Regions in the $(m_{Z_2}, M_{\tilde{x}})$ plane which satisfy all conditions discussed in the text. $\mu = 700$ GeV is assumed. Values used for $M_{\tilde{W}}$ and θ are shown in each figure.

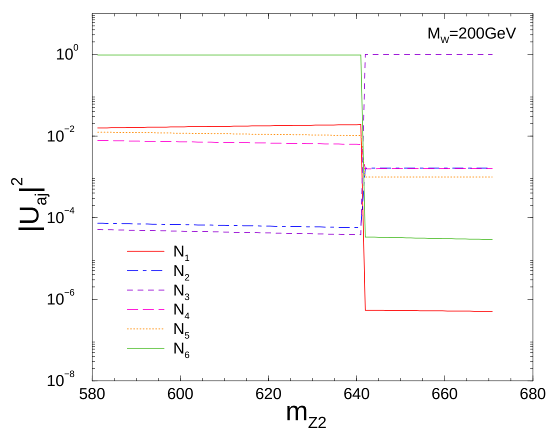

Fig. 4 Composition of the lightest neutralino in the case of $\theta = -0.253$, $M_{\tilde{x}} = 4.5$ TeV and $M_{\tilde{W}} = 200$ GeV.

which can give the lower bounds of m_{Z_2}. On the other hand, the boundary expressed by a vertical line at a larger m_{Z_2} is caused by the mass bounds of the lightest neutral Higgs scalar given in (27). Since we fix a value of μ, larger values of u make values of λ smaller. As we can see from (13) and also as discussed in the previous part, smaller values of λ can decrease the upper bounds of the Higgs mass and make it below the current experimental bounds at small values of $\tan\beta$. These features confine the mass of the new gauge boson into narrow regions as shown in each figure. Solutions for other values of μ have the similar features to those shown in this figure except that the solutions for a larger μ value move to the region with larger values of m_{Z_2}.

From Figs. 3 we find that all solutions are found in the regions where the lightest

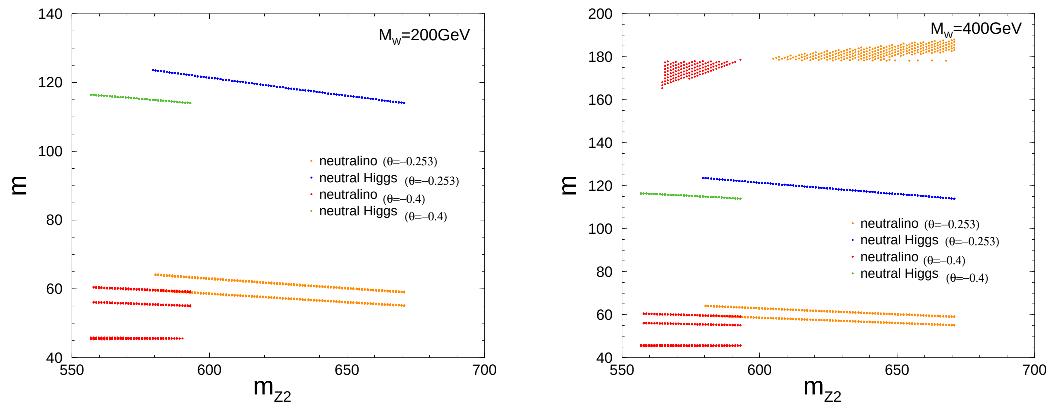

Fig. 5 Masses of the lightest neutral Higgs and the lightest neutralino in the cases of $\theta = -0.4$ and -0.253. $M_{\tilde{W}}$ is shown in each figure.

neutralino is dominated by the singlino. In order to see the change of the composition of the lightest neutralino at the lower boundary of the regions of singlino domination, in Fig. 4 we plot values of $|U_{\ell j}|^2$ defined in eq. (9) for m_{Z_2} in the case of $\theta = -0.253$, $M_{\tilde{x}} = 4.5$ TeV, and $M_{\tilde{W}} = 200$ GeV. It shows that the singlino domination of the lightest neutralino suddenly turns into the bino domination within the narrow region of m_{Z_2}. The other cases with different values of $M_{\tilde{W}}$ and μ also show the similar behavior to this case on the same boundary. For larger values of M_W than 400 GeV we find the allowed regions in the same places in the $(m_{Z_2}, M_{\tilde{x}})$ plane since the composition of the lightest neutralino is almost fixed there.

In Fig. 5 we plot the masses of the lightest neutralino and the lightest neutral Higgs scalar obtained from the solutions shown in Fig. 3. Since larger values of $M_{\tilde{x}}$ produce smaller masses for the singlino dominated lightest neutralino as discussed in section 2, we can find the correspondence of the solutions in Fig. 3 to the mass $m_{\tilde{\chi}_\ell^0}$ of the the lightest neutralino shown in Fig. 5. We can find from Fig. 5 that there are three types of possibilities for the annihilation of the lightest neutralinos to explain the CDM abundance. They are characterized by the lightest neutralino mass. The first possibility is characterized by the relatively small mass such as ~ 45 GeV. It appears commonly in the all figures in case of $\theta = -0.4$. However, we cannot find this type of solutions in case of $\theta = -0.253$. In these solutions the effective annihilation of the lightest neutralinos seems to be induced through a Z_1 exchange. It can realize the appropriate CDM abundance only if $m_{\tilde{\chi}_\ell^0} \sim m_{Z_1}/2$ and the Higgsino components are suitably contained in the singlino

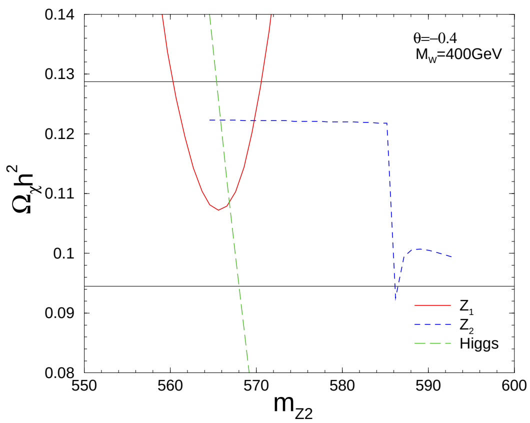

Fig. 6 Dominant processes which realize the suitable relic abundance $\Omega_\chi h^2$ in the case of $\theta = -0.4$ and and $M_{\tilde{W}} = 400$ GeV. In order to realize three types of solution $M_{\tilde{x}}$ is chosen such as 10.3 TeV for the Z_1 exchange, 7.7 TeV for the Higgs exchange, and 1.9 TeV for the Z_2 exchange.

dominated lightest neutralino. This conditions seems to be satisfied only in the restricted values of θ such as $-0.46 \lesssim \theta \lesssim -0.34$ as found in the left-hand figure of Fig. 2. The second one is characterized by the mass tuned into the narrow region around $m_{h^0}/2$. These solutions appear as a result of the Higgs pole enhancement of the annihilation. The third one is characterized by the large neutralino masses such as $\sim 170-180$ GeV. In this case the annihilation is considered to be mainly induced by a Z_2 exchange. The large mass of the singlino dominated neutralino can enhance the factor b_f in eq. (20) to realize the appropriate abundance. This type of solutions appear only for large values of $M_{\tilde{W}}$ such as 400 GeV or more. This is because the large $M_{\tilde{W}}$ is necessary for the heavy singlino dominated neutralino to be the lightest one.

In order to check these points, we adopt the solutions for $\theta = -0.4$ and $M_W = 400$ GeV in Fig. 3 and calculate $\Omega_\chi h^2$ by taking account of only the annihilation process discussed above for each case. We choose the values of $M_{\tilde{x}}$ as 10.3 TeV, 7.7 TeV and 1.9 TeV, where the three types of solutions appear. We plot the results in Fig. 6. This figure justifies the above discussion. Dominant components of the lightest neutralino in each case are found to be

$$
\begin{array}{llll}
Z_1 \text{ exchange } (M_{\tilde{x}} = 10.3 \text{ TeV}): & \mathcal{N}_6 \sim 97.6\%, & \mathcal{N}_5 \sim 1.4\%, & \\
h^0 \text{ exchange } (M_{\tilde{x}} = 7.7 \text{ TeV}): & \mathcal{N}_6 \sim 97.4\%, & \mathcal{N}_5 \sim 1.4\%, & \\
Z_2 \text{ exchange } (M_{\tilde{x}} = 1.9 \text{ TeV}): & \mathcal{N}_6 \sim 91\%, & \mathcal{N}_1 \sim 7\%, & \mathcal{N}_5 \sim 1\%.
\end{array}
$$

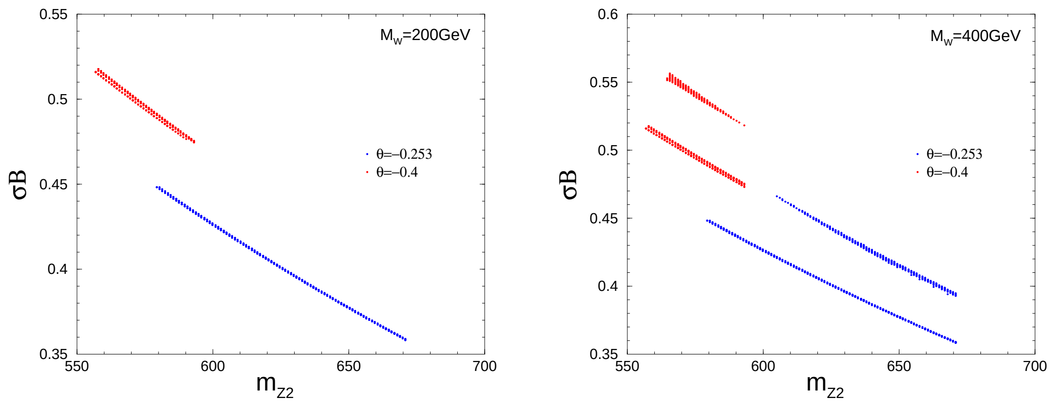

Fig. 7 An expected cross section for the Z_2 decay into the dilepton final states at the LHC with $\sqrt{s} = 14$ TeV in the cases of $\theta = -0.4$ and -0.253. $\mu = 700$ GeV is assumed and $M_{\tilde{W}}$ is shown in each figure.

For the heavy neutralino solutions the lightest neutralino can be heavier than the ordinary Z_1. Then we may need to take account of additional new final states such as W^+W^- and $Z_1 h^0$ which can be mediated by the Z_2 exchange. However, since these processes are generally suppressed in comparison with the ones with final states $f\bar{f}$ in the present models, these effects seem to be safely neglected in the annihilation of the singlino dominated lightest neutralino. Thus, the present results are considered to be a good approximation even in the case of $m_{\tilde{\chi}_\ell^0} > m_{Z_1}$.

In Fig. 7 we plot the predicted cross section σB of the Z_2 decay into the dilepton final states $(e^+e^-, \mu^+\mu^-)$ at the LHC with $\sqrt{s} = 14$ TeV for the models defined by $\theta = -0.4$ and -0.253. This figure also shows that the LHC can easily find Z_2 in the present models. The models defined by the different values of θ can also be distinguished from each other. In the previous part we discussed that the dilepton channel may be suppressed compared with the case where the lightest neutralino is composed of only the MSSM contents, since the Z_2 decay into the neutralino pairs can be enhanced in the case of the singlino dominated lightest neutralino. This may be generally expected if $\hat{S}$ has the larger charge of $\mathrm{U}(1)_x$ compared with the charged leptons. In the present models this is the case and the aspect is shown in the figure for $M_{\tilde{W}} = 400$ GeV in Fig. 7. In the case of $M_{\tilde{W}} = 400$ GeV the appearance of new branches can be understood from this reasoning. They correspond to the solutions with the heavier lightest neutralino. Thus the Z_2 decay to the neutralino sector is suppressed and the cross section to the dilepton final states is enhanced.

4 Summary

We have studied the phenomenology of the models with an extra U(1) which can give an elegant weak scale solution for the μ problem and have the singlino dominated lightest neutralino. Such models can be constructed by introducing an SM singlet chiral superfield $\hat{S}$ to the MSSM. The vacuum expectation value u of the scalar component of $\hat{S}$ generates the μ term and also breaks the extra U(1) symmetry. The singlino dominated lightest neutralino is realized by assuming the extra U(1) gaugino to be heavier than other gauginos. Under this assumption, the lightest neutralino can be dominated by the fermionic component of $\hat{S}$ even if the vacuum expectation value u is large enough to make the extra U(1) gauge field Z_2 satisfy the current experimental constraints. We studied the case with $\tan\beta = \sqrt{Q_1/Q_2}$ where the ZZ' mixing constraint is automatically satisfied even for rather small value of u. Although this fact is known, the study of that case seems not to be done in the Z' phenomenology.

The neutralino sector of the models is different from the currently known similar type models. In the NMSSM and the nMSSM the singlino dominated lightest neutralino usually appears for the small values of u. In the models with an extra U(1) in which the masses of the gauginos are assumed to be the same order, the lightest neutralino is expected to be very similar to that of the MSSM because of a large value of u. On the other hand, in the present models the singlino dominated lightest neutralino can appear in a very different situation from these and this can make their phenomenology distinguishable from others.

We have focussed our study on the extra U(1) models which are derived from E_6. And we have discussed typical phenomenological consequences in the cases where the singlino dominated lightest neutralino has the suitable relic density as a CDM candidate. We have shown that there are non-negligible parameter regions which satisfy the currently known phenomenological constraints as long as the model satisfies $\tan\beta \simeq \sqrt{Q_1/Q_2}$. The CDM candidate has been shown to have different natures from that in both the MSSM and other type of singlet extensions of the MSSM. They may be examined in future collider experiments. We have also predicted the masses of the lightest neutral Higgs scalar, the new neutral gauge boson and the lightest neutralino. The predicted cross section for the Z_2 decay into the dilepton pairs shows that the Z_2 boson in the models can be easily found at the LHC. Searches of the neutral Higgs scalar and the Z_2 boson at the LHC are expected to give us useful informations for the models with an extra U(1) which may give the solution for the μ problem.

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Appendix

In this appendix we present the interaction Lagrangian relevant to possible decay processes of Z_2 and also formulas for the decay width for them [14].

(i)$Z_2 \to f\bar{f}$

Interaction terms relevant to this decay process are

$$\mathcal{L} = \frac{g_x}{2}\bar{f}\gamma_\mu(q_v^f - q_a^f\gamma_5)fZ_2^\mu,$$
$$q_v^f = \frac{1}{2}\left(Q(f_L) + Q(f_R)\right), \qquad q_a^f = \frac{1}{2}\left(Q(f_L) - Q(f_R)\right). \tag{30}$$

The decay width can be calculated as

$$\Gamma_{f\bar{f}} = c_f \frac{g_x^2}{4}\frac{m_{Z_2}}{12\pi}\left[q_v^{f2}\left(1 + 2\frac{m_f^2}{m_{Z_2}^2}\right) + q_a^{f2}\left(1 - 4\frac{m_f^2}{m_{Z_2}^2}\right)\right]\sqrt{1 - 4\frac{m_f^2}{m_{Z_2}^2}}, \tag{31}$$

where c_f equals to 1 for leptons and 3 for quarks. Because of various reasons, the extra fermions in the present models can be considered to be sufficiently heavy and then neglected in the final states for the Z_2 decay. In that case it is justified to confine f into ordinary quarks and leptons.

(ii) $Z_2 \to \tilde{f}_{L,R}\tilde{f}^*_{L,R}$

Mass matrices of the sfermions have a similar structure to that in the MSSM except that there are additional contributions from the extra U(1) D-term. It can play a crucial role since it may give large negative contributions depending on the models. They can be written as

$$\mathcal{M}_{\tilde{f}}^2 = \begin{pmatrix} M^2_{\tilde{f}_L\tilde{f}_L} & M^2_{\tilde{f}_L\tilde{f}_R} \\ M^{2\dagger}_{\tilde{f}_L\tilde{f}_R} & M^2_{\tilde{f}_R\tilde{f}_R} \end{pmatrix}, \tag{32}$$

where each component of this matrix is a 3×3 matrix with respect to flavors. However, if flavor mixing is small enough, each flavor can be treated independently and $\mathcal{M}^2_{\tilde{f}}$ can be reduced into a 3×3 matrix. In that case the component of $\mathcal{M}^2_{\tilde{f}}$ can be expressed as

$$M^2_{\tilde{f}_L\tilde{f}_L} = \tilde{m}_0^2 + |m_f|^2 + (T_3 - Q_f s_W^2)m_Z^2\cos 2\beta + Q_x(f_L)\tilde{m}_D^2,$$
$$M^2_{\tilde{f}_R\tilde{f}_R} = \tilde{m}_0^2 + |m_f|^2 + Q_f s_W^2 m_Z^2\cos 2\beta + Q_x(f_R)\tilde{m}_D^2,$$
$$M^2_{\tilde{f}_L\tilde{f}_R} = -m_f^*(A^* - \lambda u R_f), \tag{33}$$

where (T_3, R_f) takes $(\frac{1}{2}, \cot\beta)$ and $(-\frac{1}{2}, \tan\beta)$ for an up- and down-sector of squarks and sleptons, respectively. m_f and Q_f are the mass and the electric charge of a fermion f.

We assume the universality of soft scalar masses $\tilde{m}_f^2$ and A parameters. The D-term contribution of the extra U(1) has the expression

$$\tilde{m}_D^2 = \frac{1}{2} g_x^2 \left(Q_1 v_1^2 + Q_2 v_2^2 + Q_S u^2 \right). \qquad (34)$$

We should note that this contribution may be negative to realize the light sfermions depending on the value of θ introduced in eq. (24) which defines the extra U(1) models.

Interaction terms relevant to this decay process are given by

$$\mathcal{L} = \frac{g_x}{2} (q_v^f \pm q_a^f) \tilde{f}_{L,R}^* i \overleftrightarrow{\partial_\mu} f_{L,R} Z_2^\mu, \qquad (35)$$

where v_f and a_f are defined in eq.(30). The decay width of this process can be written as

$$\Gamma_{\tilde{f}_{L,R} \tilde{f}_{L,R}^*} = c_f \frac{g_x^2}{4} \frac{m_{Z_2}}{48\pi} (q_v^f \pm q_a^f)^2 \left(1 - 4 \frac{m_{\tilde{f}_{L,R}}^2}{m_{Z_2}^2} \right)^{3/2}. \qquad (36)$$

In this derivation the LR mixing $M_{\tilde{f}_L \tilde{f}_R}^2$ in eq. (32) is neglected. This is considered to be a good approximation except for the stop sector.

(iii) $Z_2 \to H^+ H^-$

The charged Higgs sector has the same structure as that of the MSSM except that there are additional contributions to the mass eigenvalue generated by the coupling $\lambda \hat{S} \hat{H}_1 \hat{H}_2$ as in the NMSSM. The mass eigenvalue and its eigenstate can be expressed as

$$m_{H^\pm}^2 = m_W^2 \left(1 - \frac{2\lambda^2}{g_2^2} \right) + \frac{2A\lambda u}{\sin 2\beta}, \qquad H^\pm = H_1^\pm \sin\beta + H_2^\pm \cos\beta. \qquad (37)$$

The interaction Lagrangian relevant to this process is given by

$$\mathcal{L} = \frac{g_x}{2} (Q_1 \sin^2\beta - Q_2 \cos^2\beta) H^+ i \overleftrightarrow{\partial_\mu} H^- Z_2^\mu, \qquad (38)$$

where $G^\pm$ are would-be Goldstone bosons. Using this interaction Lagrangian, we can derive the decay width for this process as

$$\Gamma_{H^+H^-} = \frac{g_x^2}{4} \frac{m_{Z_2}}{48\pi} (Q_1 \sin^2\beta - Q_2 \cos^2\beta)^2 \left(1 - 4 \frac{m_{H^\pm}^2}{m_{Z_2}^2} \right)^{3/2}. \qquad (39)$$

(iv) $Z_2 \to W^\pm H^\mp$

The relevant interaction Lagrangian for this process is

$$\mathcal{L} = \frac{g_x}{2} (Q_1 + Q_2) \sin\beta \cos\beta \left[(m_W H^- W_\mu^+ + H^- i \overleftrightarrow{\partial_\mu} G^+) Z_2^\mu + \text{h.c.} \right]. \qquad (40)$$

Using this Lagrangian, the decay width for this process can be calculated as

$$\begin{aligned}\Gamma_{W^\pm H^\mp} &= \frac{g_x^2}{4}\frac{m_{Z_2}}{48\pi}(Q_1+Q_2)^2\sin^2\beta\cos^2\beta\left[1-2\frac{m_W^2+m_{H^\pm}^2}{m_{Z_2}^2}+\frac{(m_W^2-m_{H^\pm}^2)^2}{m_{Z_2}^4}\right]^{1/2}\\ &\times \left[1+2\frac{5m_W^2-m_{H^\pm}^2}{m_{Z_2}^2}+\frac{(m_W^2-m_{H^\pm}^2)^2}{m_{Z_2}^4}\right]. \qquad (41)\end{aligned}$$

(v) $Z_2 \to Z_1\phi_\alpha$

The relevant interaction Lagrangian is written as

$$\mathcal{L}=\sum_{\alpha=1}^{3}2\frac{g_x}{2}m_Z(Q_1\mathcal{O}_{1\alpha}\cos\beta-Q_2\mathcal{O}_{2\alpha}\sin\beta)(Z_\mu+\partial_\mu G^0)\phi_\alpha Z_2^\mu. \qquad (42)$$

where we denote a would-be Goldstone boson as G^0. The matrix $\mathcal{O}$ and the mass eigenvalues $m_{\phi_\alpha}^2$ are defined in eq. (12) and also in the text below it. The decay width for this process can be derived as

$$\begin{aligned}\Gamma_{Z\phi_\alpha} &= \sum_{\alpha=1}^{3}\frac{g_x^2}{4}\frac{m_{Z_2}}{48\pi}(Q_1\mathcal{O}_{1\alpha}\cos\beta-Q_2\mathcal{O}_{2\alpha}\sin\beta)^2\left[1-2\frac{m_{\phi_\alpha}^2+m_{Z_1}^2}{m_{Z_2}^2}+\frac{(m_{\phi_\alpha}^2-m_{Z_1}^2)^2}{m_{Z_2}^4}\right]^{1/2}\\ &\times \left[1+2\frac{5m_{Z_1}^2-m_{\phi_\alpha}^2}{m_{Z_2}^2}+\frac{(m_{\phi_\alpha}^2-m_{Z_1}^2)^2}{m_{Z_2}^4}\right]. \qquad (43)\end{aligned}$$

(vi) $Z_2 \to \phi_\alpha P_A$

P_A is the CP odd Higgs scalar. The relevant interaction Lagrangian for this process is written as

$$\mathcal{L}=\sum_{\alpha=1}^{3}\frac{g_x}{2}\left(Q_1\mathcal{O}_{1\alpha}\frac{u\sin\beta}{N}+Q_2\mathcal{O}_{2\alpha}\frac{u\cos\beta}{N}+Q_S\mathcal{O}_{3\alpha}\frac{v\sin\beta\cos\beta}{N}\right)Z_2^\mu P_A\overleftrightarrow{\partial_\mu}\phi_\alpha, \qquad (44)$$

where N is defined in eq. (16) and in the text below it. Using this, we can derive the decay width as

$$\begin{aligned}\Gamma_{\phi_\alpha P_A} &= \frac{g_x^2}{4}\frac{m_{Z_2}}{48\pi}\frac{1}{N^2}(Q_1\mathcal{O}_{1\alpha}u\sin\beta+Q_2\mathcal{O}_{2\alpha}u\cos\beta+Q_S\mathcal{O}_{3\alpha}v\sin\beta\cos\beta)^2\\ &\times \left[1-2\frac{m_{\phi_\alpha}^2+m_{P_A}^2}{m_{Z_2}^2}+\frac{(m_{\phi_\alpha}^2-m_{P_A}^2)^2}{m_{Z_2}^4}\right]^{3/2}. \qquad (45)\end{aligned}$$

(vii) $Z_2 \to \tilde\chi_a^0\tilde\chi_b^0$

The neutralino sector is discussed in detail in the text. The interaction Lagrangian for the neutralinos $\tilde\chi_a^0$ and Z_2 is given by

$$\mathcal{L}=\sum_{a,b=1}^{6}g_{ab}\bar{\tilde\chi}_a^0\gamma_\mu\gamma_5\tilde\chi_b^0 Z_2^\mu, \qquad (46)$$

where an effective coupling g_{ab} is defined as

$$g_{ab} = \sum_{j=4}^{6} \frac{g_x}{2} Q(\mathcal{N}_j) U_{aj} U^*_{bj}. \tag{47}$$

The definition of the matrix U is given in eq. (9). Using this interaction Lagrangian, the decay width is derived as

$$\begin{aligned}
\Gamma_{\tilde{\chi}^0_a \tilde{\chi}^0_b} &= \frac{g_{ab}^2 m_{Z_2}}{12\pi} \left(1 - \frac{1}{2}\delta_{ab}\right) \left[1 - \frac{m_a^2 + m_b^2}{2m_{Z_2}^2} - \frac{(m_a^2 - m_b^2)^2}{2m_{Z_2}^4} - \frac{3m_a m_b}{m_{Z_2}^2}\right] \\
&\times \left[1 - 2\frac{m_a^2 + m_b^2}{m_{Z_2}^2} + \frac{(m_a^2 - m_b^2)^2}{m_{Z_2}^4}\right]^{1/2},
\end{aligned} \tag{48}$$

where m_a stands for the mass eigenvalue of the neutralino $\tilde{\chi}^0_a$.

(viii) $Z_2 \to \tilde{C}^+_a \tilde{C}^-_b$

The chargino sector has the same structure as that of the MSSM. The mass terms can be expressed as

$$\frac{1}{2}(-i\tilde{\lambda}_{W^-}, \tilde{H}_1^-) \begin{pmatrix} M_{\tilde{W}} & \sqrt{2} m_W \sin\beta \\ \sqrt{2} m_W \cos\beta & -\lambda u \end{pmatrix} \begin{pmatrix} -i\tilde{\lambda}_{W^+} \\ \tilde{H}_2^+ \end{pmatrix} + \text{h.c.}. \tag{49}$$

If we define the mass eigenstates as

$$\begin{aligned}
\begin{pmatrix} \psi_1^+ \\ \psi_2^+ \end{pmatrix} &= V_+ \begin{pmatrix} -i\tilde{\lambda}_{W^+} \\ \tilde{H}_2^+ \end{pmatrix}, \qquad V_+ \equiv \begin{pmatrix} \cos\phi_+ & -\sin\phi_+ \\ \sin\phi_+ & \cos\phi_+ \end{pmatrix}, \\
\begin{pmatrix} \chi_2^- \\ \chi_1^- \end{pmatrix} &= V_- \begin{pmatrix} -i\tilde{\lambda}_{W^-} \\ \tilde{H}_1^- \end{pmatrix}, \qquad V_- \equiv \begin{pmatrix} \cos\phi_- & -\sin\phi_- \\ \sin\phi_- & \cos\phi_- \end{pmatrix},
\end{aligned} \tag{50}$$

the mixing angles $\phi_\pm$ have analytic expressions

$$\begin{aligned}
\tan 2\phi_+ &= \frac{2\sqrt{2} m_W (M_{\tilde{W}} \sin\beta - \lambda u \cos\beta)}{2m_W^2(\sin^2\beta - \cos^2\beta) + \lambda^2 u^2 - M_{\tilde{W}}^2}, \\
\tan 2\phi_- &= \frac{2\sqrt{2} m_W (M_{\tilde{W}} \cos\beta - \lambda u \sin\beta)}{2m_W^2(\cos^2\beta - \sin^2\beta) + \lambda^2 u^2 - M_{\tilde{W}}^2}.
\end{aligned} \tag{51}$$

Using this basis, the mass terms can be transformed into

$$\frac{1}{2}(\chi_1^-, \chi_2^-) \begin{pmatrix} m_1 & 0 \\ 0 & m_2 \end{pmatrix} \begin{pmatrix} \psi_1^+ \\ \psi_2^+ \end{pmatrix} + \text{h.c.}, \tag{52}$$

where the mass eigenvalues are represented as

$$\begin{aligned} m_1 &= (M_{\tilde{W}} \sin\phi_- + \sqrt{2} m_W \cos\beta \cos\phi_-) \cos\phi_+ \\ &\qquad -(\sqrt{2} m_W \sin\beta \sin\phi_- - \lambda u \cos\phi_-) \sin\phi_+, \\ m_2 &= (M_{\tilde{W}} \cos\phi_- - \sqrt{2} m_W \cos\beta \sin\phi_-) \sin\phi_+ \\ &\qquad +(\sqrt{2} m_W \sin\beta \cos\phi_- + \lambda u \sin\phi_-) \cos\phi_+. \end{aligned} \tag{53}$$

Since only the Higgsinos $\tilde{H}_{1,2}$ interact with Z_2, the interaction Lagrangian relevant to the Z_2 decay is given by

$$\mathcal{L} = \frac{g_x}{2} \sum_{a,b=1}^{2} \bar{\tilde{C}}_a \gamma_\mu (v_{ab} + a_{ab}\gamma_5) \tilde{C}_b Z_2^\mu, \tag{54}$$

where $\tilde{C}_a^{\pm T} = (\psi_a^\pm, \chi_a^\pm)$, and effective couplings v_{ab} and a_{ab} are defined as

$$\begin{aligned} v_{11} &= -\frac{Q_1}{2}\sin^2\phi_- + \frac{Q_2}{2}\sin^2\phi_+, \qquad a_{11} = \frac{Q_1}{2}\sin^2\phi_- + \frac{Q_2}{2}\sin^2\phi_+, \\ v_{22} &= -\frac{Q_1}{2}\cos^2\phi_- + \frac{Q_2}{2}\cos^2\phi_+, \qquad a_{22} = \frac{Q_1}{2}\cos^2\phi_- + \frac{Q_2}{2}\cos^2\phi_+, \\ v_{12} &= v_{21} = \frac{Q_1}{2}\sin\phi_-\cos\phi_- - \frac{Q_2}{2}\sin\phi_+\cos\phi_+, \\ a_{12} &= a_{21} = -\frac{Q_1}{2}\sin\phi_-\cos\phi_- - \frac{Q_2}{2}\sin\phi_+\cos\phi_+. \end{aligned} \tag{55}$$

Using these couplings, the decay width for this process is derived as

$$\begin{aligned} \Gamma_{\tilde{C}_a\tilde{C}_b} &= \frac{g_x^2}{4}\frac{m_{Z_2}}{48\pi}\left[(v_{ab}^2 + a_{ab}^2)\left(1 - \frac{m_a^2 + m_b^2}{2m_{Z_2}^2} - \frac{(m_a^2 - m_b^2)^2}{2m_{Z_2}^4}\right) + 3(v_{ab}^2 - a_{ab}^2)\frac{m_a m_b}{m_{Z_2}^2}\right] \\ &\times \left[1 - 2\frac{m_a^2 + m_b^2}{m_{Z_2}^2} + \frac{(m_a^2 - m_b^2)^2}{m_{Z_2}^4}\right]^{1/2}. \end{aligned} \tag{56}$$

Finally, the branching ratio $B(Z_2 \to XY)$ used in the text is defined by

$$B(Z_2 \to XY) = \frac{\Gamma(Z_2 \to XY)}{\Gamma_{\rm tot}}, \tag{57}$$

where the total decay width $\Gamma_{\rm tot}$ of Z_2 is

$$\begin{aligned} \Gamma_{\rm tot} &= \sum_f \Gamma_{f\bar{f}} + \sum_{\tilde{f}} \Gamma_{\tilde{f}_{L,R}\tilde{f}^*_{L,R}} + \Gamma_{H^+H^-} + \Gamma_{W^\pm H^\mp} \\ &+ \sum_{i=1}^{3}\Gamma_{Z\phi_i} + \sum_{j=1}^{3}\Gamma_{\phi_j P_A} + \sum_{a,b=1}^{6}\Gamma_{\tilde{N}_a\tilde{N}_b} + \sum_{a,b=1}^{2}\Gamma_{\tilde{C}_a\tilde{C}_b}. \end{aligned} \tag{58}$$

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